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Unifying the PST and the auxiliary tensor field formulations of 4*D* self-duality

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A B S T R A C T

We unify the Lorentz- and $O(2)$ duality-covariant approach to 4*D* self-dual theories by Pasti, Sorokin and Tonin (PST) with the formulation involving an auxiliary tensor field. We present the basic features of the new hybrid approach, including symmetries of the relevant generalized PST action. Its salient peculiarity is the unique form of the realization of the PST gauge symmetries. The corresponding transformations do not affect the auxiliary tensor field, which guarantees the self-duality of the nonlinear actions in which the $O(2)$ invariant interactions are constructed out of the tensor field.

1. Introduction

Self-duality is one of the central concepts of gauge theories and string theory. The notorious examples of self-dual 4*D* systems are the renowned Born–Infeld theory and other duality-invariant models of nonlinear electrodynamics. Reconciling the manifest symmetry under duality rotations [1–3] with the manifest Lorentz invariance becomes possible in the formulations with auxiliary fields [4–8]. The most economic approach requires just a single scalar auxiliary field entering the action non-polynomially. This formulation was originally developed for the free self-dual tensor fields [4,5,8] and, later on, was extended to nonlinear models of branes and their coupling to supergravity actions [9–13]. Introducing the interaction into the self-dual covariant actions is a non-trivial task as the interaction terms should satisfy a consistency condition generalizing that of [14].

On the other hand, there exists a universal approach to duality-invariant 4*D* theories based on employing the auxiliary tensor (bispinor) fields [15–17]. It came out as a by-product of studying $\mathcal{N}=3$, 4*D* Born–Infeld theory in the harmonic superspace formulation [18]. Within this approach, the interaction part of the action is constructed solely out of the auxiliary tensor fields and is manifestly duality-invariant. Though the whole action is not duality-invariant, on shell it leads to the equations of motion which, together with the Bianchi identity, are covariant under the duality rotations. After elimination of the auxiliary fields by their equations of motion, the resulting system automatically obeys the general nonlinear Gaillard–Zumino consistency condition [14]. The whole set of the self-dual actions of nonlinear 4*D* electrodynamics thus proves to be in the one-to-one correspondence with the appropriate auxiliary interactions. In Refs. [15–17] various nonlinear self-dual models were explicitly constructed in this way. The supersymmetric versions of the approach with the auxiliary tensor fields were worked out in [19].

The aim of the present Letter is to elaborate on a new formulation of the self-dual 4*D* actions with tensor fields, such that they enjoy the manifest duality invariance off shell. This goal is pursued by properly extending the construction of Pasti–Sorokin–Tonin (PST) [4,5,8]. The striking feature of the hybrid formulation is that the renowned PST gauge symmetry transformations have a universal form, irrespective of the precise structure of the self-interaction.

The Letter is organized as follows. In Section 2 we recall the structure and the symmetries of the original PST action of the free duality-symmetric Maxwell field in 4*D*. In Section 3 we extend the PST formulation by introducing an auxiliary tensor field and discuss the duality invariance, as well as the symmetry structure of the proposed action. It allows a direct generalization to the interaction case, without affecting the form of the PST gauge transformations. Section 4 contains a brief discussion of the relations of the new action to the previously known non-covariant duality-symmetric actions of the 4*D* Maxwell field. Summary and conclusions are collected in the final part of the Letter.

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2. PST action and its symmetries

We start with a brief discussion of the standard PST approach to $4D$ self-dual theories within the original second order formulation of [4,5,8]. The following action[1]

$$S_{PST} = \frac{1}{2}\int d^4x\big[-v^m \tilde{F}^a_{mn}\delta_{ab}\tilde{F}^{bnl}v_l + v^m F^b_{mn}\epsilon_{ab}\tilde{F}^{anl}v_l\big] \tag{1}$$

describes the dynamics of the duality-symmetric Maxwell field A^a_m, $a = 1,2$, with the field strength $F^a_{mn} = \partial_m A^a_n - \partial_n A^a_m$. Its dual is defined by

$$\tilde{F}^a_{mn} = \frac{1}{2}\epsilon_{mnrl}F^{arl}. \tag{2}$$

The other entities entering the action (1) are the PST scalar $a(x)$ in the specific non-polynomial combination v_m

$$v_m = \frac{\partial_m a(x)}{\sqrt{(\partial a)^2}} \tag{3}$$

and the $O(2)$ invariant tensors δ_{ab} and ϵ_{ab} ($\epsilon_{12} = 1$, $\epsilon_{21} = -1$). The role of the PST scalar is to make the Lorentz covariance of the action manifest; the $O(2)$ tensors ensure the manifest invariance of the action under the $O(2)$ duality rotations of the vector fields A^a_m, $\delta A^a_m = \omega\epsilon_{ab}A^b_m$.

The general variation of the action (1) is calculated to be

$$\begin{aligned}\delta S_{PST} &= -\frac{1}{2}\int d^4x\big[\delta F^a_{mn} - \delta v_m\big(v\cdot\mathcal{F}^a\big)_n\big]\epsilon^{mnrl}v_r\epsilon_{ab}\big(v\cdot\mathcal{F}^b\big)_l \\ &= -\int d^4x\bigg(\delta A^a_m - \frac{\delta a}{\sqrt{(\partial a)^2}}\big(v\cdot\mathcal{F}^a\big)_m\bigg) \\ &\quad\times\epsilon^{mnrl}\partial_n\big(v_r\epsilon_{ab}\big(v\cdot\mathcal{F}^b\big)_l\big),\end{aligned} \tag{4}$$

where we have omitted the total derivative term and introduced

$$\mathcal{F}^a_{mn} = F^a_{mn} + \epsilon_{ab}\tilde{F}^b_{mn}, \tag{5}$$

with $(v\cdot\mathcal{F}^a)_n := v^m\mathcal{F}^a_{mn}$. As follows from (4), the action (1), besides the invariance under $U(1)$ local gauge transformations of the Maxwell fields A^a_m, reveals the invariance under the following special gauge symmetries (the so-called PST symmetries [4,5,8]):

$$\delta_I A^a_m = \partial_m a(x)\Phi_a(x), \qquad \delta_I a(x) = 0, \tag{6}$$

$$\delta_{II} a(x) = \varphi(x), \qquad \delta_{II} A^a_m = \frac{\varphi(x)}{\sqrt{(\partial a)^2}}\big(v\cdot\mathcal{F}^a\big)_m. \tag{7}$$

These two PST symmetries (below we refer to them as PST-I and PST-II) play different roles. The PST-I symmetry is needed to reduce, by fixing the gauge parameters $\Phi_a(x)$, the vector field equation of motion

$$\epsilon^{mnrl}\partial_n\big[v_r\epsilon_{ab}\big(v\cdot\mathcal{F}^b\big)_l\big] = 0 \tag{8}$$

to the self-duality condition[2]

$$\mathcal{F}^a_{mn} = F^a_{mn} + \epsilon_{ab}\tilde{F}^b_{mn} = 0. \tag{9}$$

The second PST symmetry (PST-II) guarantees that the presence of the PST scalar $a(x)$, which is needed for the manifest Lorentz covariance of the PST action, does not increase the number of the initial degrees of freedom: This field can be completely gauged away. Indeed, its equation of motion

$$\partial_m\bigg[\frac{1}{\sqrt{(\partial a)^2}}\epsilon^{mnrl}v_n\epsilon_{ab}\big(v\cdot\mathcal{F}^a\big)_r\big(v\cdot\mathcal{F}^b\big)_l\bigg] = 0 \tag{10}$$

does not contain any additional information and is trivially satisfied on shell. However, a gauge fixing of the PST scalar breaks the manifest Lorentz covariance of the model, resulting in the manifestly duality-invariant but non-covariant formulation of Schwarz and Sen [2] (see Section 4 for details).

3. Manifestly covariant self-dual action with tensor auxiliary fields

As shown in [15–17], introducing the auxiliary tensor fields allows one to drastically simplify the problem of finding the $O(2)$ duality-invariant interactions. It was of obvious interest to generalize the auxiliary tensor field formulation in such a way that the $O(2)$ duality becomes the manifest off-shell symmetry of the total action.

This goal motivated us to consider the following modification of the PST action (1)

$$\begin{aligned}S = \int d^4x\,\mathcal{L}'_{PST} \equiv \int d^4x\big[&\mathcal{L}_{PST} + v^m\tilde{F}^a_{mn}\delta_{ab}\tilde{F}^{bnl}v_l \\ &+ v^m V_{mn}V^{nr}v_r + v^m\tilde{V}_{mn}\tilde{V}^{nr}v_r + 2v^m V_{mn}\tilde{F}^{2nl}v_l \\ &- 2v^m\tilde{V}_{mn}\tilde{F}^{1nl}v_l\big].\end{aligned} \tag{11}$$

Integrating out an unconstrained auxiliary field V_{mn} takes as back to the original PST action. Indeed, the equations of motion for V_{mn} read

$$v_{[m}\big(V_{n]p} + \tilde{F}^2_{n]p}\big)v^p - \frac{1}{2}v^t\epsilon_{tsmn}\big(\tilde{V}^{sp} - \tilde{F}^{1sp}\big)v_p = 0, \tag{12}$$

whence

$$v^m V_{mn} = -v^m\tilde{F}^2_{mn}, \qquad v^m\tilde{V}_{mn} = v^m\tilde{F}^1_{mn}. \tag{13}$$

Substitution of these expressions into (11) leaves us with $\mathcal{L}_{PST}$ as the "on-shell" Lagrangian. Note that the relations (13) in fact enable to express the whole V_{mn} in terms of v_m and the field strengths F^a_{mn} (the number of the independent relations in (13) is just 6, and V_{mn} has 6 independent components). It is convenient to present the corresponding expressions in the bispinor notation, using the definitions

$$\begin{aligned}&V_{mn} - i\tilde{V}_{mn} = (\tilde{\sigma}_m\sigma_n)^{\dot\alpha}_{\dot\beta}\bar{V}^{\dot\beta}_{\dot\alpha}, \qquad V_{mn} + i\tilde{V}_{mn} = -(\sigma_m\tilde{\sigma}_n)^{\alpha}_{\beta}V^{\beta}_{\alpha}, \\ &v_m = \frac{1}{2}(\tilde{\sigma}_m)^{\dot\alpha\beta}v_{\beta\dot\alpha} = \frac{1}{2}(\sigma_m)_{\beta\dot\alpha}\tilde{v}^{\dot\alpha\beta}, \qquad v_{\beta\dot\alpha}\tilde{v}^{\dot\alpha\rho} = \delta^{\rho}_{\beta},\end{aligned} \tag{14}$$

and the similar ones for F^a_{mn}. Then Eqs. (13) amount to the set

$$\begin{aligned}&v_{\beta\dot\xi}V^{\beta}_{\xi} + v_{\xi\dot\alpha}\bar{V}^{\dot\alpha}_{\dot\xi} = i\big[v_{\beta\dot\xi}\big(F^2\big)^{\beta}_{\xi} - v_{\xi\dot\alpha}\big(\bar{F}^2\big)^{\dot\alpha}_{\dot\xi}\big], \\ &v_{\beta\dot\xi}V^{\beta}_{\xi} - v_{\xi\dot\alpha}\bar{V}^{\dot\alpha}_{\dot\xi} = v_{\beta\dot\xi}\big(F^1\big)^{\beta}_{\xi} - v_{\xi\dot\alpha}\big(\bar{F}^1\big)^{\dot\alpha}_{\dot\xi},\end{aligned} \tag{15}$$

and hence we find

$$\begin{aligned}&V^{\alpha}_{\xi} =: I^{\alpha}_{\xi}(F) = \frac{1}{2}\big(F^1 - iF^2\big)^{\alpha}_{\xi} - \frac{1}{2}v_{\xi\dot\beta}\tilde{v}^{\dot\rho\alpha}\big(\bar{F}^1 - i\bar{F}^2\big)^{\dot\beta}_{\dot\rho}, \\ &\bar{V}^{\dot\alpha}_{\dot\xi} = \overline{\big(V^{\alpha}_{\xi}\big)}.\end{aligned} \tag{16}$$

This solution could be equivalently derived in a more direct way, starting from the action (11) in the bispinor notation.

The extended action (11), like its pure PST prototype, is $O(2)$ duality invariant. Indeed, the involved quantities possess the following $O(2)$ transformation properties:

$$\begin{aligned}&\delta F^a_{mn} = \omega\epsilon_{ab}F^b_{mn}, \\ &\delta V_{mn} = \omega\tilde{V}_{mn} \quad\Leftrightarrow\quad \delta\tilde{V}_{mn} = -\omega V_{mn}.\end{aligned} \tag{17}$$

[1] We use the conventions $g_{mn} = \mathrm{diag}(1,-1,-1,-1)$, $\epsilon_{0123} = 1$.

[2] Details of this procedure may be found, e.g., in [20], or, more recently, in [23].

One can join $V_{mn}, \tilde{V}_{mn}$ into the doublet $V^a_{mn} := (V_{mn}, \tilde{V}_{mn})$, $\tilde{V}^a_{mn} = \epsilon_{ab} V^b_{mn}$, and rewrite the V-dependent terms in (11) in the manifestly $O(2)$ invariant form

$$v^m V^a_{mn} \delta_{ab} V^{bnr} v_r + 2 v^m V^a_{mn} \epsilon_{ab} \tilde{F}^{bnl} v_l. \tag{18}$$

We observe that for preserving the manifest duality invariance it is enough to add a *single* auxiliary tensor field, still keeping a double set of the gauge potentials.

Thus by construction the action (11) is manifestly Lorentz- and duality-invariant. However, like in the original PST approach [4,5,8], the covariance of the action is ensured by the PST scalar $a(x)$, the auxiliary nature of which is guaranteed by the PST symmetries. Therefore, we are led to find how (if any) the transformation laws (6), (7) are modified upon introducing the auxiliary tensor field V_{mn}.

The straightforward computations lead to the following expression for the general variation of (11) (modulo a total derivative):

$$\begin{aligned}\delta S = &-\int d^4x \left(\delta A^a_m - \frac{\delta a}{\sqrt{(\partial a)^2}} \left(v \cdot \widehat{\mathcal{F}}^a\right)_m\right) \epsilon^{mnrl} \partial_n \left(v_r \epsilon_{ab} \left(v \cdot \widehat{\mathcal{F}}^b\right)_l\right) \\ &+ 2 \int d^4x\, \delta V_{mn} \left(v^m \left[-V^{nr} - \tilde{F}^{2nr}\right] v_r \right. \\ &\left. + \frac{1}{2} v_s \epsilon^{sumn} \left[\tilde{V}_{ur} - \tilde{F}^1_{ur}\right] v^r \right),\end{aligned} \tag{19}$$

where

$$\widehat{\mathcal{F}}^1_{mn} := F^1_{mn} - \tilde{F}^2_{mn} - 2V_{mn}, \qquad \widehat{\mathcal{F}}^2_{mn} = F^2_{mn} + \tilde{F}^1_{mn} - 2\tilde{V}_{mn}. \tag{20}$$

The variation (19) vanishes under the standard local $U(1)$ transformations of the gauge fields A^a_m, as well as under the following modified PST-type transformations

$$\delta_I A^a_m = \partial_m a(x) \Phi_a(x), \qquad \delta_I a(x) = 0, \qquad \delta_I V_{mn} = 0, \tag{21}$$

$$\begin{aligned}&\delta_{II} a(x) = \varphi(x), \qquad \delta_{II} A^a_m = \frac{\varphi(x)}{\sqrt{(\partial a)^2}} \left(v \cdot \widehat{\mathcal{F}}^a\right)_m, \\ &\delta_{II} V_{mn} = 0.\end{aligned} \tag{22}$$

We see that only the PST-II transformations are actually modified. It is very important to note that the modified PST transformations (21), (22) do not affect the auxiliary tensor field V_{mn}. This peculiarity has a great impact on the structure of admissible interaction terms.

Just due to this notable property, in the nonlinear case we can add, to the bilinear action (11), an arbitrary auxiliary interaction $\mathcal{E}(A) = \frac{1}{2}A + O(A^2)$, where A is the quartic $O(2)$ invariant variable

$$A = \left(\operatorname{Tr} V^2\right)\left(\operatorname{Tr} \bar{V}^2\right) = \frac{1}{16}\left[\left(V^{mn} V_{mn}\right)^2 + \left(V^{mn} \tilde{V}_{mn}\right)^2\right], \tag{23}$$

with $(\operatorname{Tr} V^2) = V^\beta_\alpha V^\alpha_\beta = \overline{(\operatorname{Tr} \bar{V}^2)}$. The free solution $V^\alpha_\beta = I^\alpha_\beta(F)$ (16) can be easily generalized to the interaction case

$$V^\beta_\alpha = I^\beta_\alpha(F) - (v \bar{V} \bar{v})^\beta_\alpha \left(\operatorname{Tr} V^2\right) \mathcal{E}_A, \tag{24}$$

where $\mathcal{E}_A = d\mathcal{E}/dA$. The resulting action

$$S_{int} = \int d^4x \left[\mathcal{L}'_{PST} + \mathcal{E}(A)\right] \tag{25}$$

preserves all symmetries of the bilinear action, including the gauge PST symmetries. Thus it describes some self-dual systems for any choice of the $O(2)$ invariant interaction. After solving Eq. (24) (e.g., by recursions), one is left with the highly nonlinear action in terms of the field strengths F^a_{mn} and the auxiliary scalar field $a(x)$.[3] The PST-II transformations also become nonlinear. It is remarkable that, before eliminating V_{mn}, the PST transformations have the universal form (21), (22).

4. Non-covariant self-dual action with tensor auxiliary field

Let us briefly discuss the relation of the covariant action (11) to the previously proposed duality-invariant actions of [2,21,16].

The PST-II symmetry of the action (22) makes it possible to fix the gauge $v_m = \delta^0_m$. As a result, the covariant action (11) turns into the following duality-symmetric non-covariant action:

$$S_{n.c.} = \int d^4x \left[\frac{1}{2} B^a_k \delta_{ab} B^b_k + \frac{1}{2} B^a_k \epsilon_{ab} E^b_k + 2 V^a_k \epsilon_{ab} B^b_k + V^a_k \delta_{ab} V^b_k\right]. \tag{26}$$

Here, we have introduced $V_{0i} = V^1_i = V_i$, $\tilde{V}_{0i} = V^2_i = U_i$.[4] The action (26) is a non-covariant gauge-fixed "magnetic" version of the covariant action proposed in [21] (see Appendix C of [16] for details of deriving (26)[5]).

After integrating out the auxiliary fields V^a_k, the action (26) takes the form

$$S_{SS} = \int d^4x \left[-\frac{1}{2} B^a_k \delta_{ab} B^b_k + \frac{1}{2} B^a_k \epsilon_{ab} E^b_k\right], \tag{27}$$

which is none other than the Schwarz–Sen non-covariant duality-invariant action [2]. This result is of course not surprising, because the elimination of the tensor auxiliary field in the gauge-unfixed action (11) takes the latter just back to the PST action. On the other hand, the PST action (1) is a covariantization of the Schwarz–Sen action (27).

Extending the action (26) to the non-trivial interaction comes about as follows

$$\begin{aligned}L_{n.c.} = &\frac{1}{2} B^a_k \delta_{ab} B^b_k + \frac{1}{2} B^a_k \epsilon_{ab} E^b_k + 2 V^a_k \epsilon_{ab} B^b_k \\ &+ V^a_k \delta_{ab} V^b_k + \mathcal{E}(A),\end{aligned} \tag{28}$$

where $\mathcal{E}(A)$ is the same function of the manifestly $O(2)$ duality-invariant variable (23) as in the covariant action (25). In the $3D$ notation it is constructed out of the $3D$ components of V_{mn} as

$$A = \frac{1}{4}(U_k U_k)^2 + \frac{1}{4}(V_k V_k)^2 - \frac{1}{2}(U_k U_k)(V_i V_i) + (V_i U_i)^2. \tag{29}$$

The corresponding auxiliary equations are

$$\begin{aligned}&V^1_k + B^2_k + \frac{1}{2} \frac{\partial A}{\partial V^1_k} \mathcal{E}_A = 0, \qquad \mathcal{E}_A \equiv \frac{\partial \mathcal{E}(A)}{\partial A}, \\ &V^2_k - B^1_k + \frac{1}{2} \frac{\partial A}{\partial V^2_k} \mathcal{E}_A = 0.\end{aligned} \tag{30}$$

The general non-covariant action (28) can be obtained as a gauge-fixed version of the general covariant action (25) which enjoys both PST gauge symmetries.

[3] The $O(2)$ invariant interactions with derivatives of V_{mn} are also admissible, leading to self-dual actions with derivatives on the gauge field strengths [17].

[4] $4D$ indices are split into the $1+3$ set as $m = (0, i)$. We denote $F^a_{0i} = E^a_i$, $F^a_{ij} = \epsilon_{ijk} B^a_k$, $\tilde{F}^a_{0i} = B^a_i$; in our notation $\epsilon^{0123} = -1$, so $\epsilon^{0ijk} = -\epsilon_{ijk}$ and $\epsilon_{0ijk} = \epsilon_{ijk}$.

[5] In fact, in [16] an "electric" version of (26) was derived, using the trick suggested in [21]; the action (26) follows from the "electric" version through a discrete duality transformation.

5. Conclusions

To summarize, we have proposed the new approach to the $4D$ self-dual nonlinear electrodynamics systems, which is a symbiosis of the Pasti–Sorokin–Tonin covariant duality-invariant approach [4,5,8] and the approach of [15–17] involving auxiliary tensor fields. The new approach inherits the advantages of both approaches just mentioned. On the one hand, it is manifestly Lorentz and $O(2)$ duality covariant. On the other hand, it provides a simple way of constructing self-dual actions with a non-trivial interaction. We have studied the symmetry structure of the proposed action and established its relation to the duality-symmetric approaches of [2,21]. The most sound feature of the action constructed is the universal form of the gauge PST transformations off shell, before eliminating the auxiliary tensor fields by their equations of motion. These transformations do not affect the auxiliary fields at all, the feature that makes it possible to construct invariant interactions from these fields without breaking any symmetry of the free action. Note that PST actions with additional auxiliary fields were considered before (see, e.g., [8,22]), however the approach we follow here is entirely different, since it is not related to any dualization of the PST scalar field.

An important feature of our formulation is that it ensures a consistent way of adding a non-trivial interaction to the free actions, with the guarantee that the emerging nonlinear system is self-dual. Recently, the general structure of nonlinear interacting self-dual actions within the PST approach was analyzed in [23]. It was found there that the invariance of the whole action under the PST-type transformations amounts to the fundamental consistency condition of the Gaillard–Zumino type [14]. It would be interesting to establish the precise links of our approach with this general analysis.

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